A Field-Theoretic Parametrization of Low-Energy Nucleon Form Factors

Brian D. Serot*

Department of Physics and Nuclear Theory Center
Indiana University, Bloomington, IN 47405

(Dated: December 27, 2009)

Abstract

A field-theoretic parametrization is proposed for nucleon electromagnetic form factors at momentum transfer less than 600 MeV. The parametrization is part of a larger effective field theory lagrangian that is Lorentz covariant and chiral symmetric, and that has been used to successfully describe bulk and single-particle properties of medium to heavy mass nuclei. The parametrization is based on vector meson dominance and a derivative expansion of nucleon couplings to the electromagnetic fields. At lowest order in the expansion, it is possible to fit all four parameters to modern data on the rms radii of the nucleon form factors. At next-to-leading order it is possible to fit the form factors to within a few percent up to momentum transfers of 600 MeV. The vector meson dominance contributions are crucial in this fit, since a simple expansion in powers of momentum transfer would require many, many terms to achieve comparable accuracy. The ability to fit single-nucleon form factors up to 600 MeV momentum transfer makes possible the study of two-body electromagnetic exchange currents within this effective field theory framework.

PACS numbers: 14.20.Dh; 25.30.Bf; 12.40.Vv; 11.10.-z

*Electronic address: serot@indiana.edu

I. INTRODUCTION

Lorentz-covariant meson–baryon effective field theories of the nuclear many-body problem (often called *quantum hadrodynamics* or QHD) have been known for many years to provide a realistic description of the bulk properties of nuclear matter and heavy nuclei. (For reviews, see Refs. [1, 2, 3, 4, 5, 6].) Recently, a QHD effective field theory (EFT) has been proposed [7, 8, 9, 10, 11, 12] that includes all the relevant symmetries of the underlying QCD. In particular, the spontaneously broken $SU(2)_L \times SU(2)_R$ chiral symmetry is realized nonlinearly. The motivation for this EFT and some calculated results are discussed in Refs. [6, 7, 13, 14, 15, 16, 17, 18, 19, 20, 21].

This QHD EFT has three desirable features: (1) It uses the same degrees of freedom to describe the currents and the strong-interaction dynamics; (2) It respects the same internal symmetries, both discrete and continuous, as the underlying QCD (before and after electromagnetic interactions are included); and (3) Its parameters can be calibrated using strong-interaction phenomena, like πN scattering and the empirical properties of finite nuclei (as opposed to electroweak interactions with nuclei). It thus provides a natural framework, based on a single lagrangian, for discussing the roles of one-body and two-body currents in nuclear electromagnetic interactions.

The nucleon electromagnetic (EM) structure (form factors) is described in this EFT using a combination of vector meson dominance (VMD) [7, 22, 23, 24, 25, 26, 27] and a derivative expansion for nucleon interactions with the EM field. In the applications of this EFT to nuclear structure noted above, however, only the lowest-order derivative couplings were included, so that the form factors provided an accurate description of the single-nucleon electron scattering data only up to roughly 250 MeV momentum transfer. In contrast, if one is to study *two-body* (exchange) currents, one must reproduce the single-nucleon form factors accurately up to at least 600 MeV momentum transfer, where two-body contributions are expected to be visible. This momentum scale should be accessible in this low-energy hadronic EFT [7, 8, 28].

Our motivation for this study of nucleon form factors is twofold. First, we want to update the lowest-order fits of Ref. [7] to include the large amount of low-energy, high-precision data that became available in the early 2000's. Second, we extend the fits to the next order in momentum transfer and show that the form factors will accurately reproduce the empirical results up to roughly 600 MeV momentum transfer. This will make them suitable for studies of exchange currents within the QHD EFT.

In the past ten or fifteen years, much new data on the nucleon EM form factors have been obtained using both unpolarized electron scattering and polarization transfer. (For recent reviews, see Refs. [29, 30].) There have also been numerous attempts at fitting the improved data set; for example, see Refs. [31, 32, 33]. For the present study, we are most interested in the work of Kelly [32], who achieved excellent fits with a small number of free parameters. In particular, the fits are good enough over the momentum transfer range of interest to us that we will simply fit our EFT parameters to Kelly's analytic results rather than to the data itself. Since our best fits reproduce Kelly's at the few percent level, this procedure is justified.

One of our interesting results is that a straightforward Q^2 expansion of Kelly's analytic results is inadequate unless many, many terms are retained. (Here $Q^2 \equiv -q^2$ is the square of the spacelike four-momentum transfer.) The presence of the VMD contributions in the EFT approach greatly improves the situation. Moreover, it is important to include the new

EFT parameters in such a way that the error is minimized for the whole relevant range of Q^2, not just $Q^2 \to 0$.

This paper is organized as follows: In Sec. II, we return to the lowest-order parametrization of Ref. [7] and re-fit the EFT parameters to the new data set. This allows us to determine mean-square radii for all four form factors (neutron/proton–electric/magnetic). In Sec. III, we extend the EFT lagrangian by introducing new parameters and use them to fit the higher-momentum transfer behavior of the form factors. Sec. IV is a brief Summary.

II. RE-FIT OF LOWEST-ORDER PARAMETERS

In this section, we consider the form factors as described in Ref. [7]. We follow the conventions of Refs. [7, 12]. Rather than work with the Dirac (F_1) and Pauli (F_2) form factors, defined in terms of the nucleon EM vertex as

$$\Gamma^\mu = F_1(Q^2)\gamma^\mu + F_2(Q^2)\frac{i\sigma^{\mu\nu}q_\nu}{2M} \ , \tag{1}$$

where M is the nucleon mass and F_2 contains the anomalous magnetic moment, here we will primarily concentrate on the Sachs form factors

$$G_E(Q^2) = F_1(Q^2) - \tau F_2(Q^2) \ , \quad G_M(Q^2) = F_1(Q^2) + F_2(Q^2) \ , \tag{2}$$

where $\tau \equiv Q^2/4M^2 \equiv -q^2/4M^2$ in terms of the four-momentum q^μ, and we have not distinguished the charge states. The charge states are written in terms of the isoscalar (0) and isovector (1) parts as, for example,

$$F_p = F^{(0)} + F^{(1)} \ , \quad F_n = F^{(0)} - F^{(1)} \ . \tag{3}$$

The simple parametrizations used by Kelly [32] take the form

$$G(Q^2) = \frac{\sum_{k=0}^{n} a_k \tau^k}{1 + \sum_{k=1}^{n+2} b_k \tau^k} \ , \tag{4}$$

which guarantees the correct asymptotic dependence at large Q^2: $G(Q^2) \propto Q^{-4}$. This will not concern us, as we are interested in parametrizing the form factors at small Q^2. With $n = 1$ and $a_0 = 1$, this parametrization gives excellent fits to G_{Ep}, G_{Mp}/μ_p, and G_{Mn}/μ_n (where μ_i is the full magnetic moment) using four parameters each [32]. For G_{En}, Kelly follows the so-called Galster parametrization [34]:

$$G_{En}(Q^2) = \frac{A\tau}{1 + B\tau}\, G_D(Q^2) \ , \tag{5}$$

where the dipole form factor is

$$G_D(Q^2) \equiv \frac{1}{(1 + Q^2/\Lambda^2)^2} \ , \quad \Lambda^2 = 0.71\,\mathrm{GeV}^2 \ , \tag{6}$$

and A and B are fitted parameters.

For our parametrization, we use set Q2 of Ref. [7]. This provides an accurate fit to bulk and single-particle nuclear properties and leaves the anomalous coupling to the isoscalar

TABLE I: Coupling parameters from set Q2 [7]. Note that the nucleon couplings to the omega and rho mesons (g_v and g_ρ) are determined from the empirical properties of nuclei.

$\beta^{(0)}$	$\beta^{(1)}$	g_v	g_ρ	f_v	f_ρ
0.01181	−0.1847	12.2148	8.5572	0	4.264

vector meson ("omega") undetermined; we will determine it here. We will need the mass parameters

$$\begin{aligned} M &= 939\,\mathrm{MeV} = 4.7585\,\mathrm{fm}^{-1} \ , \\ m_v &= 782\,\mathrm{MeV} = 3.963\,\mathrm{fm}^{-1} \ , \\ m_\rho &= 770\,\mathrm{MeV} = 3.902\,\mathrm{fm}^{-1} \ , \end{aligned} \tag{7}$$

the anomalous magnetic moments

$$\lambda_p = 1.793 \ , \quad \lambda_n = -1.913 \ , \tag{8}$$

the couplings in Table I, and the electromagnetic coupling $g_\gamma = 5.0133$, which follows from the decay width $\Gamma_{\rho^0 \to e^+e^-} = 6.8\,\mathrm{keV}$.

We now proceed with the fits to the new data. The calculations use the lagrangian of Ref. [7] and are performed at tree level. Starting with the proton electric form factor,

$$\begin{aligned} G_{Ep}(Q^2) &= F_1^{(0)} + F_1^{(1)} - \frac{Q^2}{4M^2}\,(F_2^{(0)} + F_2^{(1)}) \\ &= 1 - (\beta^{(0)} + \beta^{(1)})\frac{Q^2}{2M^2} - \frac{g_v}{3g_\gamma}\left(1 - \frac{f_v Q^2}{4M^2}\right)\frac{Q^2}{Q^2 + m_v^2} \\ &\quad - \frac{g_\rho}{2g_\gamma}\left(1 - \frac{f_\rho Q^2}{4M^2}\right)\frac{Q^2}{Q^2 + m_\rho^2} - \lambda_p \frac{Q^2}{4M^2} \ . \end{aligned} \tag{9}$$

If we define the mean-square radius as

$$r_{Ep}^2 \equiv -6 \left.\frac{\mathrm{d}G_{Ep}(Q^2)}{\mathrm{d}Q^2}\right)_{Q^2=0} \ , \tag{10}$$

then

$$\begin{aligned} r_{Ep}^2 &= \frac{1}{2}\left[6\left(\frac{\beta^{(0)}}{M^2} + \frac{2g_v}{3g_\gamma m_v^2}\right) + 6\left(\frac{\beta^{(1)}}{M^2} + \frac{g_\rho}{g_\gamma m_\rho^2}\right)\right] + \frac{3\lambda_p}{2M^2} \\ &\equiv \frac{1}{2}\left(\langle r^2\rangle_{(0)1} + \langle r^2\rangle_{(1)1}\right) + \frac{3\lambda_p}{2M^2} \ , \end{aligned} \tag{11}$$

where the mean-square radii on the right-hand side are the isoscalar and isovector values for the Dirac form factor F_1. Inserting the Q2 parameters leads to

$$(r_{Ep}^2)^{1/2} = 0.862\,\mathrm{fm} \ , \tag{12}$$

which also agrees with the result in Ref. [32].

Turning now to G_{En}, we have

$$\begin{aligned} G_{En}(Q^2) &= F_1^{(0)} - F_1^{(1)} - \frac{Q^2}{4M^2}(F_2^{(0)} - F_2^{(1)}) \\ &= -(\beta^{(0)} - \beta^{(1)})\frac{Q^2}{2M^2} - \frac{g_v}{3g_\gamma}\left(1 - \frac{f_v Q^2}{4M^2}\right)\frac{Q^2}{Q^2 + m_v^2} \\ &\quad + \frac{g_\rho}{2g_\gamma}\left(1 - \frac{f_\rho Q^2}{4M^2}\right)\frac{Q^2}{Q^2 + m_\rho^2} - \lambda_n \frac{Q^2}{4M^2} \ , \end{aligned} \tag{13}$$

so that

$$\begin{aligned} r_{En}^2 &= \frac{1}{2}\left[6\left(\frac{\beta^{(0)}}{M^2} + \frac{2g_v}{3g_\gamma m_v^2}\right) - 6\left(\frac{\beta^{(1)}}{M^2} + \frac{g_\rho}{g_\gamma m_\rho^2}\right)\right] + \frac{3\lambda_n}{2M^2} \\ &= \frac{1}{2}\left(\langle r^2\rangle_{(0)1} - \langle r^2\rangle_{(1)1}\right) + \frac{3\lambda_n}{2M^2} \ . \end{aligned} \tag{14}$$

If we set $\langle r^2\rangle_{(0)1} = \langle r^2\rangle_{(1)1}$, as in Ref. [7], we then find $r_{En}^2 = -0.127\,\mathrm{fm}^2$, in significant disagreement with Kelly's value of $-0.112 \pm 0.003\,\mathrm{fm}^2$. We conclude that the new data shows that

$$\langle r^2\rangle_{(0)1} - \langle r^2\rangle_{(1)1} = 0.0294\,\mathrm{fm}^2 \ . \tag{15}$$

With this information, together with Eq. (11), we can determine two distinct radii for the Dirac form factor:

$$\langle r^2\rangle_{(0)1}^{1/2} = 0.799\,\mathrm{fm} \ , \quad \langle r^2\rangle_{(1)1}^{1/2} = 0.780\,\mathrm{fm} \ , \tag{16}$$

in contrast to the assumption made in Ref. [7].

We now consider the magnetic form factors, beginning with

$$\begin{aligned} G_{Mp}(Q^2) &= F_1^{(0)} + F_1^{(1)} + F_2^{(0)} + F_2^{(1)} \\ &= 1 - (\beta^{(0)} + \beta^{(1)})\frac{Q^2}{2M^2} - \frac{g_v(1+f_v)}{3g_\gamma}\frac{Q^2}{Q^2 + m_v^2} \\ &\quad - \frac{g_\rho(1+f_\rho)}{2g_\gamma}\frac{Q^2}{Q^2 + m_\rho^2} + \lambda_p \ . \end{aligned} \tag{17}$$

As expected, for $Q^2 \to 0$, $G_{Mp}(Q^2) \to 1 + \lambda_p = \mu_p$. Thus we normalize G_{Mp} by dividing by μ_p, leading to the mean-square radius

$$\begin{aligned} r_{Mp}^2 &= \frac{1}{2(1+\lambda_p)}\left[6\left(\frac{\beta^{(0)}}{M^2} + \frac{2g_v}{3g_\gamma m_v^2}\right) + 6\left(\frac{\beta^{(1)}}{M^2} + \frac{g_\rho}{g_\gamma m_\rho^2}\right)\right] \\ &\quad + \frac{1}{2(1+\lambda_p)}\left(\frac{4f_v g_v}{g_\gamma m_v^2} + \frac{6f_\rho g_\rho}{g_\gamma m_\rho^2}\right) \\ &= \frac{1}{2(1+\lambda_p)}\left(\langle r^2\rangle_{(0)1} + \langle r^2\rangle_{(1)1} + (\lambda_p + \lambda_n)\langle r^2\rangle_{(0)2} + (\lambda_p - \lambda_n)\langle r^2\rangle_{(1)2}\right) \\ &= 0.7191\,\mathrm{fm}^2 \ , \end{aligned} \tag{18}$$

where the numerical value is taken from Kelly. Note that the magnetic radii depend on both the Dirac and Pauli mean-square radii.

For the neutron,

$$\begin{aligned} G_{Mn}(Q^2) &= F_1^{(0)} - F_1^{(1)} + F_2^{(0)} - F_2^{(1)} \\ &= -(\beta^{(0)} - \beta^{(1)})\frac{Q^2}{2M^2} - \frac{g_v(1+f_v)}{3g_\gamma}\frac{Q^2}{Q^2+m_v^2} \\ &\quad + \frac{g_\rho(1+f_\rho)}{2g_\gamma}\frac{Q^2}{Q^2+m_\rho^2} + \lambda_n \ . \end{aligned} \tag{19}$$

The mean square radius is

$$\begin{aligned} r_{Mn}^2 &= \frac{1}{2\lambda_n}\left[6\left(\frac{\beta^{(0)}}{M^2} + \frac{2g_v}{3g_\gamma m_v^2}\right) - 6\left(\frac{\beta^{(1)}}{M^2} + \frac{g_\rho}{g_\gamma m_\rho^2}\right)\right] \\ &\quad + \frac{1}{2\lambda_n}\left(\frac{4f_v g_v}{g_\gamma m_v^2} - \frac{6f_\rho g_\rho}{g_\gamma m_\rho^2}\right) \\ &= \frac{1}{2\lambda_n}\big(\langle r^2\rangle_{(0)1} - \langle r^2\rangle_{(1)1} + (\lambda_p+\lambda_n)\langle r^2\rangle_{(0)2} - (\lambda_p-\lambda_n)\langle r^2\rangle_{(1)2}\big) \\ &= 0.8226\,\mathrm{fm}^2 \ . \end{aligned} \tag{20}$$

With the results in Eqs. (16), (18), and (20), it is now simple algebra to determine the rms radii for the Pauli form factor, with the results

$$\langle r^2\rangle_{(0)2}^{1/2} = 1.30\,\mathrm{fm} \ , \quad \langle r^2\rangle_{(1)2}^{1/2} = 0.896\,\mathrm{fm} \ . \tag{21}$$

The isovector radius is consistent with Refs. [7, 35]. At this time, we have no way to estimate errors for any of these derived radii.

Equation (20) also shows the relationships between the Dirac and Pauli rms radii and the $\mathcal{O}(Q^2)$ parameters in the EM coupling expansion (i.e., $\beta^{(0)}, \beta^{(1)}, f_v, f_\rho$). Simple inversion of these results allows us to present parameters determined from the new single-nucleon data with the nuclear couplings of set Q2:

$$\beta^{(0)} = 0.0670 \ , \quad \beta^{(1)} = -0.2402 \ , \quad f_v = -0.3279 \ , \quad f_\rho = 4.4198 \ . \tag{22}$$

III. INCLUSION OF HIGHER-ORDER COUPLINGS

So far, we have adjusted the parameters that enter at $\mathcal{O}(\tau)$ to achieve the modern values for the four rms radii. Now we want to go to higher order in τ and see what adjustments are necessary to reproduce the Sachs form factors up to $Q \approx 600\,\mathrm{MeV}$ or $\tau \approx 0.1$. We can relate these adjustments to higher-order terms in the derivative expansion of the EM lagrangian.

For the purposes of this work, we take the nucleon part of the EM lagrangian as

$$\begin{aligned} \mathcal{L}_{\mathrm{EM}} &= -eA_\mu \overline{N}\gamma^\mu \frac{1}{2}(1+\tau_3)N - \frac{e}{4M}F_{\mu\nu}\overline{N}\lambda\sigma^{\mu\nu}N - \frac{e}{2M^2}(\partial^\nu F_{\mu\nu})\overline{N}\beta\gamma^\mu N \\ &\quad - \frac{e}{4M^3}(\partial_\nu\partial^\eta F_{\mu\eta})\overline{N}\lambda'\sigma^{\mu\nu}N - \frac{e}{M^4}(\partial^2\partial^\nu F_{\mu\nu})\overline{N}\beta'\gamma^\mu N \\ &\quad - \frac{e}{2M^5}(\partial^2\partial_\nu\partial^\eta F_{\mu\eta})\overline{N}\lambda''\sigma^{\mu\nu}N + \cdots \ , \end{aligned} \tag{23}$$

where A^μ and $F^{\mu\nu}$ are the electromagnetic fields. All of the constants $\lambda, \beta, \lambda', \beta', \lambda''$ have the isospin structure

$$\begin{aligned}\lambda &= \lambda_p \frac{1}{2}(1+\tau_3) + \lambda_n \frac{1}{2}(1-\tau_3) = \frac{1}{2}(\lambda_p + \lambda_n) + \frac{1}{2}\tau_3(\lambda_p - \lambda_n) \\ &\equiv \lambda^{(0)} + \lambda^{(1)}\tau_3 \ , \quad \text{etc.} \end{aligned} \tag{24}$$

As shown in Ref. [12], the isovector parts of these constants should be modified to include pion interactions to maintain the residual chiral symmetry in the lagrangian. When applied to the nucleon form factors, however, these pions appear only in loops, which we are not considering, so we have omitted these terms.

We have already included the λ and β parameters in the preceding section. It is easy to see that the λ' constants enter the magnetic rms radii at the same order in Q^2 as the vector meson couplings f_v and f_ρ. Thus the λ' parameters are redundant in our approach and will not be considered in the sequel. Thus the four new adjustable constants at our disposal are contained in β' and λ''.

It is straightforward to work out the Feynman rules for the new vertices and to construct the tree-level contributions to the form factors. One finds that these constants enter the Dirac and Pauli form factors at $\mathcal{O}(Q^4)$. Thus they will enter the Sachs form factors at $\mathcal{O}(\tau^2)$ and $\mathcal{O}(\tau^3)$.

Our strategy is the following: We begin with the magnetic form factors, where both new terms enter at $\mathcal{O}(\tau^2)$. We numerically adjust the coefficient of this term until we get a good fit to the form factor up to $\tau = 0.1$. Here we define a good fit as one that minimizes the maximum deviation between the fit and the data (i.e., Kelly's results [32]) throughout the whole interval $0 \le \tau \le 0.1$. The result provides two constraints between the β' and λ''. We then turn to the electric form factors, where the new terms enter at both $\mathcal{O}(\tau^2)$ and $\mathcal{O}(\tau^3)$. The coefficients of both of these terms are then adjusted, consistent with the constraint, until a good fit is obtained for $0 \le \tau \le 0.1$. This procedure provides us with four numbers that can be used to determine the β' and λ''.

Denoting the new contributions to the Sachs form factors as δG, we find

$$\delta G_{Mp}(\tau) = 16(\lambda''^{(0)} + \lambda''^{(1)} - \beta'^{(0)} - \beta'^{(1)})\tau^2 \ , \tag{25}$$

$$\delta G_{Mn}(\tau) = 16(\lambda''^{(0)} - \lambda''^{(1)} - \beta'^{(0)} + \beta'^{(1)})\tau^2 \ . \tag{26}$$

Taking the best fit to the data leads to Figs. 1 and 2, and the constraints

$$\lambda''_p - \beta'_p = 2.255 \ , \quad \lambda''_n - \beta'_n = -2.786 \ . \tag{27}$$

The fits are accurate to a few percent.

For the electric form factors, the new terms are

$$\delta G_{Ep}(\tau) = -16(\beta'^{(0)} + \beta'^{(1)})\tau^2 - 64(\lambda''^{(0)} + \lambda''^{(1)})\tau^3 \ , \tag{28}$$

$$\delta G_{En}(\tau) = -16(\beta'^{(0)} - \beta'^{(1)})\tau^2 - 64(\lambda''^{(0)} - \lambda''^{(1)})\tau^3 \tag{29}$$

Taking the best fit subject to the constraints above produces Figs. 3 and 4, and the values

$$\beta'_p = -0.9438 \ , \quad \beta'_n = 0.4688 \ . \tag{30}$$

Combining these results with the constraints determines all four new parameters.

Although the relative error in G_{En} is as large as 30%, G_{En} is small in the momentum-transfer range of interest. A more relevant comparison is given in Fig. 5, where all four of the empirical form factors are shown along with fits to G_{En} and G_{Mn}.

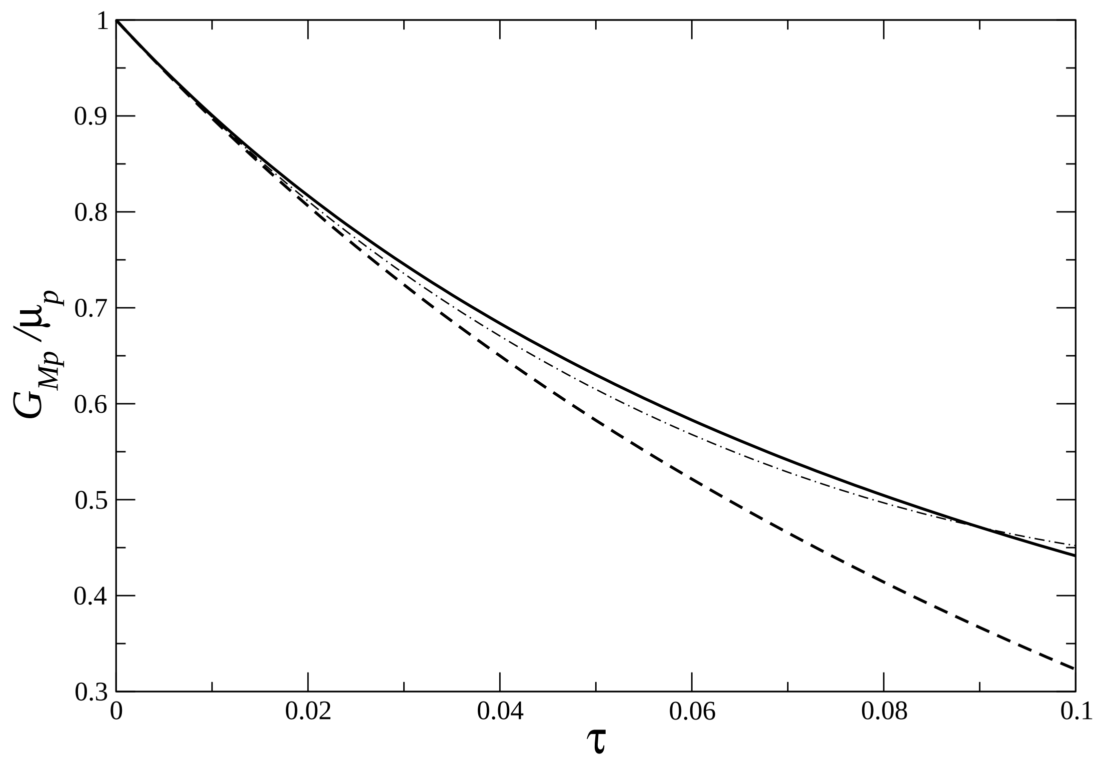

FIG. 1: Proton magnetic form factor as a function of τ. The curves represent the data (solid), the lowest-order fit (dashed), and the new fit (dot-dashed).

IV. SUMMARY

In this paper we studied a field-theoretic parametrization of single-nucleon EM form factors at low energy. This parametrization is part of a Lorentz-covariant, chiral invariant, hadronic effective field theory that was proposed to study the nuclear many-body problem. We thus have a single lagrangian that describes the nuclear structure, nuclear currents, and interaction vertices at low energies. This is a natural framework for discussing the roles of one-body and two-body currents in nuclear electromagnetic interactions.

The parametrization of the form factors is based on a combination of vector meson dominance and a derivative expansion for nucleon interactions with the EM field. At leading order in derivatives, the form factors accurately reproduce the single-nucleon electron scattering data only up to roughly 250 MeV momentum transfer. To study two-body exchange currents, however, one must reproduce the form factors accurately up to at least 600 MeV momentum transfer. This paper is proof of principle that by including the next-to-leading order (nonredundant) derivatives, one can adequately describe the form factors up to 600 MeV with our form of parametrization.

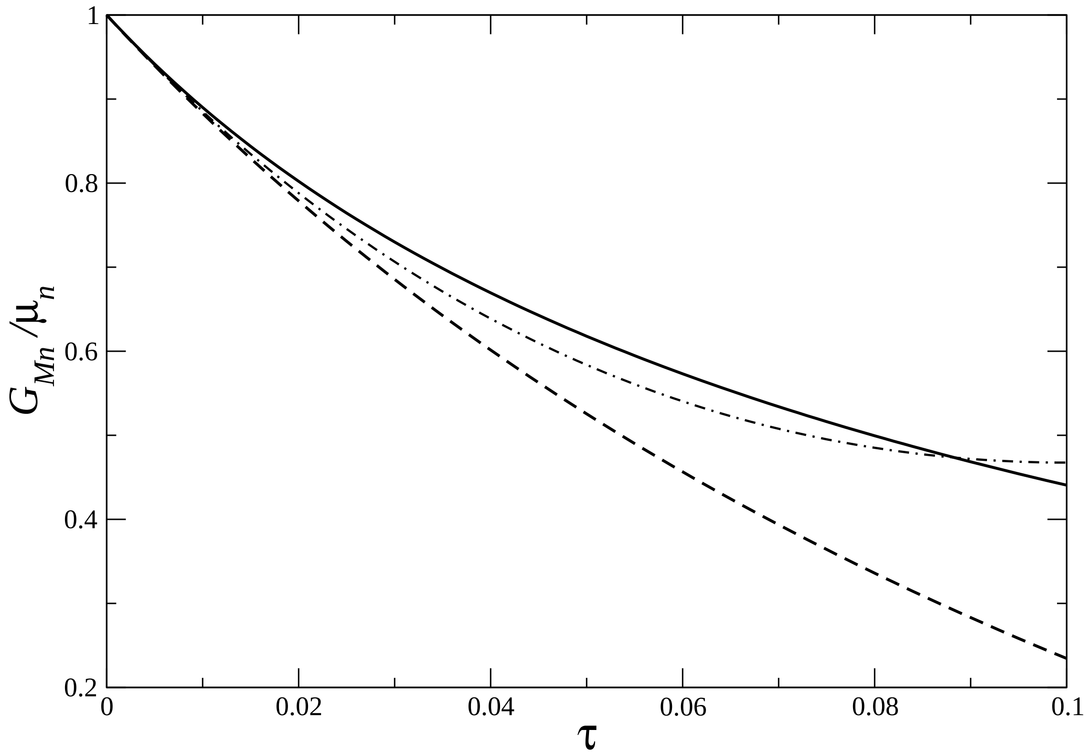

FIG. 2: Neutron magnetic form factor as a function of τ. The curves are identified as in Fig. 1.

We also revisited the leading-order parametrization and used a modern data set to evaluate the expansion coefficients. We found that it is now possible to determine all four rms radii for the Dirac and Pauli, isoscalar and isovector form factors, and thus determine all four leading-order coefficients from the single-nucleon data, unlike in Ref. [7].

It is interesting that a simple power-series expansion in powers of the momentum transfer squared would require many, many terms to reproduce the form factors up to the desired 600 MeV. (This is easy to see using ten minutes of simple calculations on MATHEMATICA with Kelly's [32] fits to the data.) Thus the vector meson dominance contributions are critical in allowing us to parametrize the desired data using only the next-to-leading order derivative terms.

While there is much work still to be done before two-body currents can be studied numerically within this effective field theory framework, we have shown that the one-body current can be adequately parametrized to make a study of two-body currents possible.

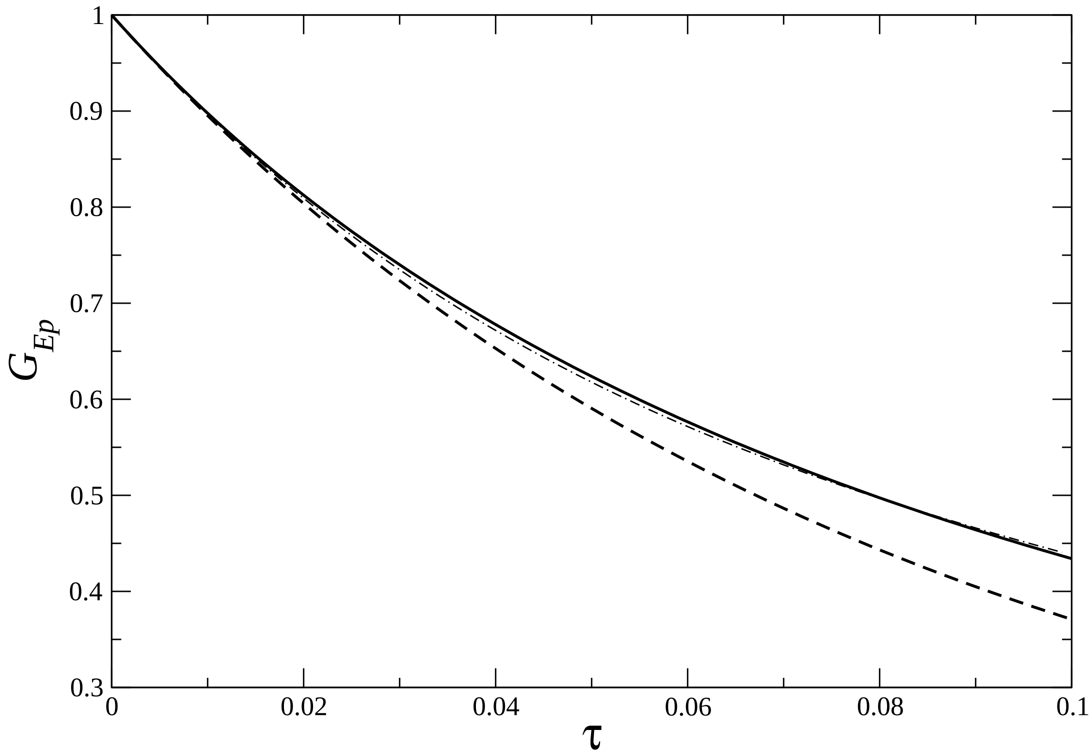

FIG. 3: Proton electric form factor as a function of τ. The curves are identified as in Fig. 1.

Acknowledgments

This work was supported in part by the Department of Energy under Contract No. DE–FG02–87ER40365.

[1] B. D. Serot and J. D. Walecka, Adv. Nucl. Phys. **16** (1986) 1.
[2] P.-G. Reinhard, Rep. Prog. Phys. **52** (1989) 439.
[3] Y. K. Gambhir, P. Ring, and A. Thimet, Ann. of Phys. **198** (1990) 132.
[4] B. D. Serot, Rep. Prog. Phys. **55** (1992) 1855.
[5] P. Ring, Prog. Part. Nucl. Phys. **37** (1996) 193.
[6] B. D. Serot and J. D. Walecka, Int. J. Mod. Phys. E **6** (1997) 515.
[7] R. J. Furnstahl, B. D. Serot, and H.-B. Tang, Nucl. Phys. **A615** (1997) 441; **A640** (1998) 505 (E).
[8] R. J. Furnstahl and B. D. Serot, Nucl. Phys. **A671** (2000) 447.
[9] R. J. Furnstahl and B. D. Serot, Nucl. Phys. **A673** (2000) 298.

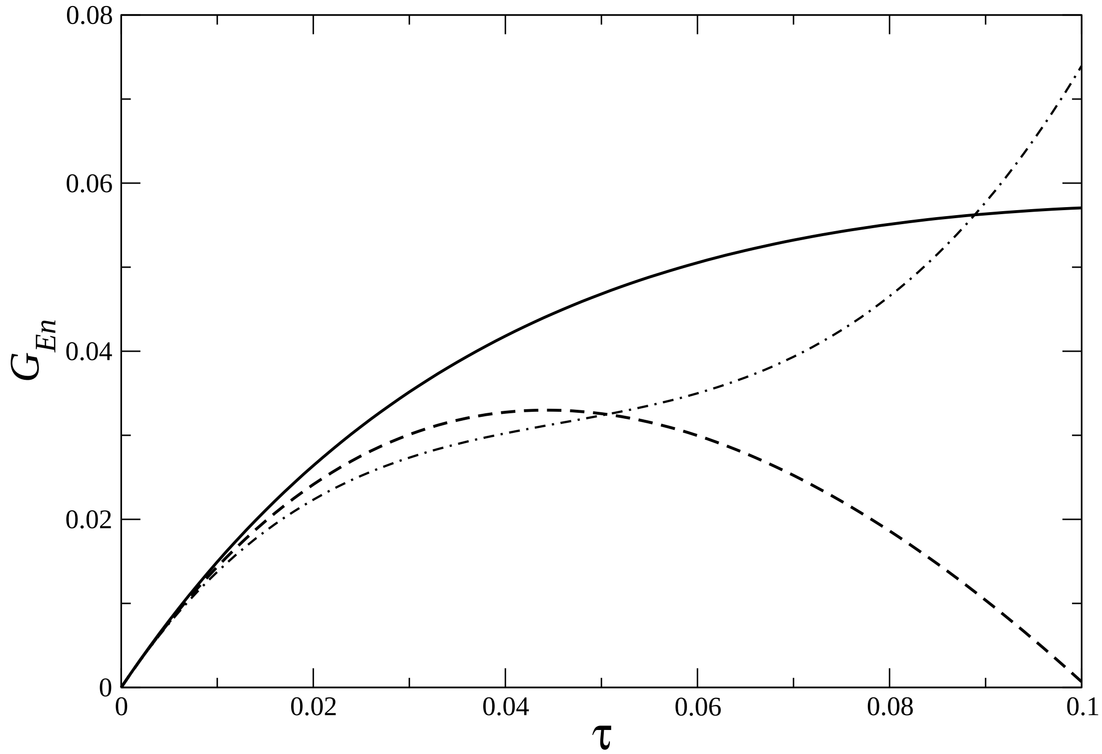

FIG. 4: Neutron electric form factor as a function of τ. The curves are identified as in Fig. 1.

[10] R. J. Furnstahl and B. D. Serot, Comments Mod. Phys. **2** (2000) A23.
[11] B. D. Serot, Lecture Notes in Physics **641**, G. A. Lalazissis, P. Ring, and D. Vretenar, eds. (Springer, Berlin Heidelberg, 2004), p. 31.
[12] B. D. Serot, Ann. of Phys. **322** (2007) 2811.
[13] M. A. Huertas, Phys. Rev. C **66** (2002) 024318; **67** (2003) 019901 (E).
[14] M. A. Huertas, Acta Phys. Polon. B **34** (2003) 4269.
[15] M. A. Huertas, Acta Phys. Polon. B **35** (2004) 837.
[16] J. McIntire, Acta Phys. Polon. B **35** (2004) 2261.
[17] J. McIntire, `arXiv:nucl-th/0507006`.
[18] J. D. Walecka, Theoretical Nuclear and Subnuclear Physics, second ed. (World Scientific, Singapore, 2004), ch. 24.
[19] J. McIntire, Y. Hu, and B. D. Serot, Nucl. Phys. **A794** (2007) 166.
[20] Y. Hu, J. McIntire, and B. D. Serot, Nucl. Phys. **A794** (2007) 187.
[21] J. McIntire, Ann. of Phys. **323** (2008) 1460.
[22] J. J. Sakurai, Ann. of Phys. **11** (1960) 1.
[23] M. Gell-Mann and F. Zachariasen, Phys. Rev. **124** (1961) 953.
[24] M. Gell-Mann, D. Sharp, and W. G. Wagner, Phys. Rev. Lett. **8** (1962) 261.

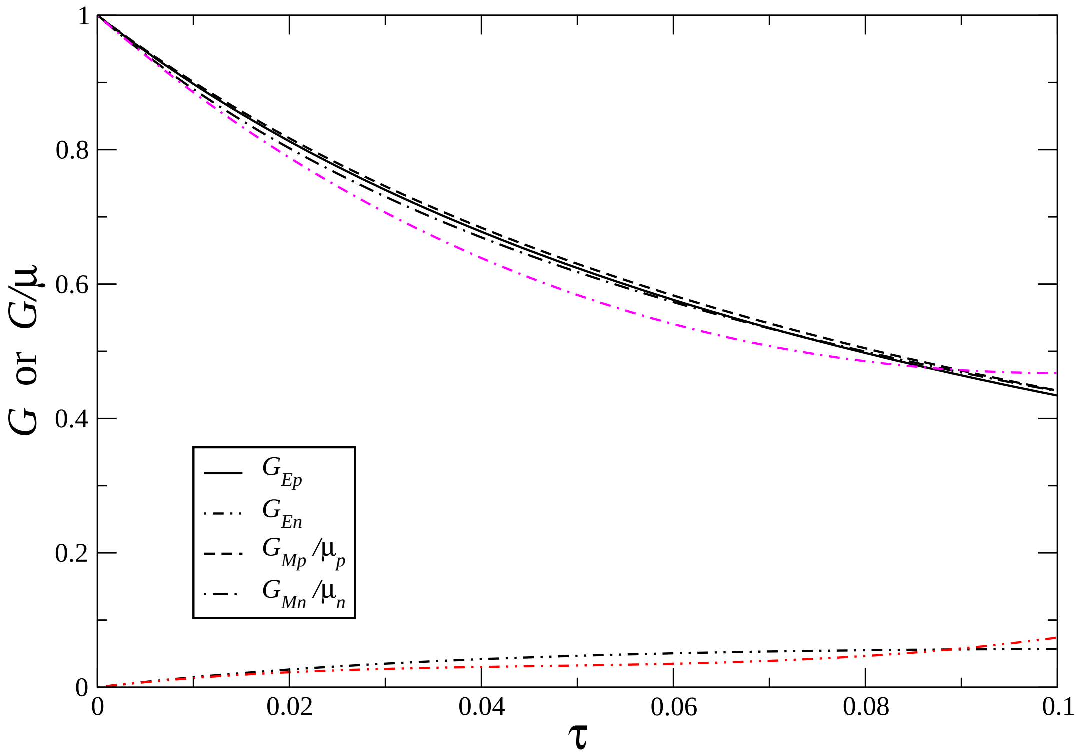

FIG. 5: All four empirical form factors [32] and fits to G_{En} and G_{Mn}.

[25] N. M. Kroll, T. D. Lee, and B. Zumino, Phys. Rev. **157** (1967) 1376.

[26] J. Bernstein, Elementary Particles and Their Currents (W. H. Freeman, San Francisco, 1968).

[27] J. J. Sakurai, Currents and Mesons (Univ. of Chicago Press, Chicago, 1969).

[28] J. J. Rusnak and R. J. Furnstahl, Nucl. Phys. **A627** (1997) 495.

[29] H. Gao, Int. J. Mod. Phys. E **12** (2003) 1.

[30] J. Arrington, C. D. Roberts, and J. M. Zanotti, J. Phys. G: Nucl. Part. Phys. **34** (2007) S23.

[31] J. J. Kelly, Phys. Rev. C **66** (2002) 065203.

[32] J. J. Kelly, Phys. Rev. C **70** (2004) 068202.

[33] W. M. Alberico, S. M. Bilenky, C. Giunti, and K. M. Graczyk, Phys. Rev. C **79** (2009) 065204.

[34] S. Galster, H. Klein, J. Moritz, K. Schmidt, D. Wegener, and J. Bleckwenn, Nucl. Phys. **B32** (1971) 221.

[35] G. E. Brown, M. Rho, and W. Weise, Nucl. Phys. **A454** (1986) 669.